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PLASMA BOUNDARY COLLISIONLESS ABSORPTION EFFECTS IN THE LOADING--ETC(U)
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PLASMA BOUNDARY COLLISIONLESS ABSORPTION EFFECTS IN THE LOADING OF RF CONDUCTORS

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ABSTRACT

This study considers the role of collisionless absorption at the boundary between an unmagnetized plasma and a conductor to which an RF potential is applied. The Landau half-space problem is examined from the energetic viewpoint to calculate the surface power absorption due ro longitudinal fields. The transverse field absorption is evaluated and the connection with the anomalous skin effect is indicated. The corresponding loading resistance is found to be significant for typical surface plasma parameters. A numerical study of pure bounce absorption heating demonstrates the physical reality of the process and uncovers the existence of a time asymptotic equilibrium between the kinetic energy of the plasma and the externally applied RF potential, independent of initial conditions or parameters.

I. Introduction

It is common practice in evaluating the loading properties of radio frequency (RF) couplers intended for plasma heating applications to consider primarily the interaction of the plasma particles with the far fields of the couplers or antennae structures. The reason for the emphasis on this particular aspect of the more general coupling problem is that in a given heating scheme one is interested in exciting collective plasma oscillations in regions distant from the location of the coupler. The loading calculations are typically aimed at evaluating the rate of energy absorption by the various physics channels of the wave damping mechanisms (collisional, collisionless, or nonlinear), and to equate such a dissipation with the actual loading which the coupling structure experiences.

In the actual experimental environment it is often found that a cold tenuous plasma is in contact with the surface of the RF couplers. These plasmas typically have densities $n_e \sim 10^{10} - 10^{12}$ cm^{-3} and electron temperatures $T_e \sim 20 - 40$ eV. Although the extent to which plasma particles interact directly with the surface of RF couplers can be minimized by the careful design of metallic guards and limiters, the fact that these particles are poorly confined, and the need to optimize the antenna coupling efficiency by placing the RF structures as close as possible to the plasma, suggest that even under controlled conditions a significant interaction may occur. Consequently, an understanding of the possible plasma boundary absorption mechanisms is desirable, since in this region there exist large electric fields of electrostatic as well as inductive origin.

The large surface electric fields can give rise to intense heating of the boundary plasma and thus cause an enhancement of the loading of the coupling structure above that calculated solely on the basis of the far field interaction. In addition, the surface heating can give rise to a variety of complicating side effects such as, change in ambipolar potential, enhanced wall sputtering, enhanced ionization of feed neutral gas, as well as multipactor breakdown between adjacent metallic surfaces.

The present study treats simple examples of collisionless plasma boundary absorption in the neighborhood of conductors to which RF potentials are applied. For simplicity, the unmagnetized plasma behavior is considered. The results obtained can be applied to magnetized plasmas in which the field lines intersect the RF coupler (e.g., whistler wave antennae, the mouth of high frequency waveguide grills, exposed couplers for ion cyclotron resonance heating), or in an approximate sense to the surface plasmas which for a variety of reasons (e.g., open field lines, enhanced transport) are poorly confined.

The prototype example of collisionless boundary absorption, i.e., the Landau half-space problem, is examined from the energetic viewpoint in Sec. II. The characteristic scaling for the power absorption due to the boundary effects is given by $|E_0|^2\bar{v}/4\pi$, where E_0 is the amplitude of the RF electric field at the boundary, and $\bar{v}$ is the average random velocity of the electrons flowing toward the boundary. This result can be simply visualized as the process of random particle flow across the low frequency ponderomotive potential due to the RF field.

The process of bounce absorption is illustrated in a numerical study presented in Sec. III. It is found that a collection of charged particles confined between two reflecting walls can be strongly heated by the application of a uniform RF field. The system exhibits a universal asymptotic state in which the average kinetic energy per particle $\langle KE \rangle /N$ is proportional to the RF potential ϕ applied between the walls, i.e., the system attains quasi-equipartition due to the stochastic collisions with the inert walls.

The absorption process associated with an evanescent inductive field is investigated in Sec. IV. Its connection with the anomalous skin effect is indicated and the corresponding loading for a RF conductor is calculated. Concluding remarks are presented in Sec. V.

II. Longitudinal Field Absorption

Consider a semi-infinite one dimensional plasma in the space $x > 0$ having uniform density n_o, and bounded by a specularly reflecting wall located at $x = 0$. It is assumed that a RF generator of frequency ω is connected to the wall, and that in the absence of plasma this geometry generates a vacuum electric field given by $\vec{E} = E_o \exp(-i\omega t)\ \hat{x}$, where $\hat{x}$ is the unit vector pointing into the plasma. This is the geometry used by Landau to study the launching and collisionless spatial damping of Langmuir waves having $\omega > \omega_p$, where ω_p is the electron plasma frequency. Here we are interested in the regime $\omega < \omega_p$, as is appropriate for low frequency RF heating schemes. Consequently, the electric field of interest is evanescent, i.e., when the plasma is present the spatial dependence of the field is given by

$$\vec{E} = (E_o/\varepsilon) \exp(-i\omega t)\ [1 - (\omega_p/\omega)^2 \exp\ (-\alpha x)]\ \hat{x} \tag{1}$$

where $\varepsilon = 1 - (\omega_p/\omega)^2$, and $\alpha = k_D|\varepsilon|^{1/2}$, with $k_D = \sqrt{3}\ \omega_p/\bar{v}$ representing the Debye wavenumber. The longitudinal evanescent field in Eq.(1) need not be identified explicitly with a particular type of RF antenna, but rather it can be thought of as a fringing field which is intrinsically generated by the manner in which the RF potential is applied to the metallic conductor. Of course, the actual extent to which the fringing field samples the plasma can be controlled by the usage of Faraday shields, as is commonly done in magnetic couplers.

We are interested in evaluating the power P absorbed per unit of antenna area A exposed to the plasma

$$\frac{P}{A} = \mathrm{Re}\ \left(\int_0^\infty dx \left\langle \vec{E} \cdot \vec{j}^* \right\rangle \right) \tag{2}$$

where Re refers to the real part, $\vec{j}$ is the electron current density oscillating at frequency ω, and $\langle\ \rangle$ denotes a time average. The power absorbed by the ions, as well as the excitation of ion acoustic waves,[2,3] is neglected in this calculation because we envision a parameter regime in which $\omega \gtrsim \omega_{pi}$, $T_e/T_i \sim 1\text{-}2$, where ω_{pi} refers to the ion plasma frequency, and T_e, T_i represent the electron and ion temperatures.

The electron current $\vec{j}$ can be obtained by integrating the linearized Vlasov equation over unperturbed trajectories. It is found that there are three different contributions to the oscillating current at a given x within the plasma. One contribution j_r is due to electrons which travel in the direction of increasing x, i.e., $v > 0$. Another contribution j_ℓ is due to electrons which travel from the interior of the plasma toward the boundary, i.e., $v < 0$. Finally, there is the contribution j_b of those electrons which arrive at x after being reflected by the boundary at $x = 0$. These contributions are

$$j_r = -\frac{e^2 n_o}{m} \int_0^\infty dv \int_0^x dx' E(x') \exp\left[i\omega(x-x')/v\right] \frac{\partial}{\partial v} f_o \tag{3}$$

$$j_\ell = -\frac{e^2 n_o}{m} \int_0^\infty dv \int_\infty^x dx' E(x') \exp\left[i\omega(x-x')/v\right] \frac{\partial}{\partial v} f_o \tag{4}$$

$$j_b = -\frac{e^2 n_o}{m} \int_0^\infty dv \int_\infty^0 dx' E(x') \exp\left[i\omega(x+x')/v\right] \frac{\partial}{\partial v} f_o \tag{5}$$

where e and m refer to the electron charge and mass.

The quantities j_r, j_ℓ, and j_b are the coefficients of terms which oscillate as exp ($-$ iωt), and have been obtained by assuming that the zero order velocity distribution function f_o (v) is symmetric in v.

The individual contributions to the absorbed power are obtained by inserting Eqs.(3)-(5) into Eq.(2) and using the spatial form of the electric field given in Eq.(1). The results take the form

$$(P/A)_s = (e^2 n_o/m)\ |E_o/\varepsilon|^2 \int_0^\infty dv\ \frac{\partial f_o}{\partial v} K_s(v,\omega) \tag{6}$$

where the index s refers to the r, ℓ, b contributions identified in Eqs.(3)-(5). The K_s are given by

$$K_r = -\ (v/\omega)^2 + (\lambda-\lambda^2/2)/[\alpha^2+(\omega/v)^2] \tag{7}$$

$$K_\ell = (\lambda-\lambda^2/2)/[\alpha^2+(\omega/v)^2] \tag{8}$$

$$K_b = -\ (v/\omega)^2 + (2\lambda)/[\alpha^2+(\omega/v)^2] + \lambda^2[\alpha^2-(\omega/v)^2]/[\alpha^2+(\omega/v)^2]^2 \tag{9}$$

where $\lambda = (\omega_p/\omega)^2$.

Using the definition given for α and summing Eq.(6) over the three contributions one arrives at a compact form for the total power absorbed by the electrons

$$\frac{P}{A} = \left(\frac{\omega_p}{\omega}\right)^2 \frac{|E_o|^2}{2\pi} \int_0^\infty dv\ v^2\ \frac{\partial f_o}{\partial v} \left[\frac{1-v^2/(3\bar{v}^2)}{1+|\varepsilon|v^2/(3\bar{v}^2)}\right]^2 \tag{10}$$

Since $f_o = f_o(v/\bar{v})$, the general scaling for the longitudinal collisionless boundary absorption mechanism is seen from Eq.(10) to be

$$\frac{P}{A} = \bar{v}\ \frac{|E_o|^2}{4\pi}\ H\left(\frac{\omega_p}{\omega}\right) \tag{11}$$

where the unitless function H gives the frequency dependence of the effect. Figure 1 exhibits this dependence for a Maxwellian velocity distribution function in the frequency interval $1.1 < (\omega_p/\omega)^2 < 200$ of interest to the evanescent near field problem. For large values of (ω_p/ω) , $H \sim (23)(\omega/\omega_p)^2$.

To obtain a feeling for the order of magnitude of the power absorbed through the longitudinal near field interaction, consider the characteristic parameters $A \sim 0.3\ m^2$, $\bar{v} \sim 10^8$cm/sec, $H \sim 10^{-1}$ from Fig. 1, and $E_0 \sim 3$ KV/cm, to find from Eq.(11) that 20 KW of RF power can be transferred to the surface of the plasma in the immediate vicinity of the RF conductor.

It should be noted that by setting $\lambda = 0$ in Eqs.(7)-(9) one obtains the power absorbed for a spatially uniform electric field, i.e., in this limit Eq.(1) yields $\vec{E} = (E_0/\epsilon)\exp(-i\omega t)\hat{x}$. It is found that even under this situation, where it is clear that transit-time effects are absent, there is net power being absorbed by the electrons. In this case the absorption mechanism is entirely due to the stochastic collisions of the plasma electrons with the boundary. Its existence in plasmas has been previously pointed out by Nekrasov and Timofeev[4], and by Baños, et al.[5] in calculations motivated by wave propagation in mirror geometry. The result presented in Eq.(10) contains both the contribution of the stochastic bounce absorption mechanism and the transit time effects due to the electrons passing through the evanescent envelope of the near field.

III. Properties of Bounce Absorption

To elucidate the reality of the bounce absorption mechanism we proceed to consider an idealized system which is heated by this process. The model system consists of a collection of charged particles which move along one dimension between two perfectly reflecting walls. An external spatially uniform electric field given by $E_0 \cos(\omega t)$ points along the direction of motion, as in the half-space problem of Sec. II. This geometry is of practical interest in connection with the multipactor breakdown effect[6] which may be encountered when a tenuous plasma is permitted to enter into open ended waveguides used in high frequency heating schemes.

Since the rate of energy transfer due to the stochastic bounce process depends linearly on the average particle flux to the wall, as seen from Eq.(11), it is suggestive that a RF induced temperature runaway may occur in the two wall problem (perfect confinement). Therefore, it is of interest to inquire what is the final state attained by such a system.

Defining the distance between the two walls to be 2L, and using the scaled variables $\tau = \omega t$, $\xi = x/L$, and $A = eE_0/m\omega^2 L$, the equation of motion is simply

$$\frac{d^2}{d\tau^2}\,\xi = -\,A\cos\tau \tag{12}$$

which is to be solved with the constraint that at $\xi = \pm 1$, $d\xi/d\tau$ reverses sign. We have solved Eq.(12) numerically for a collection of particles (typically 200 to 300) over a broad range of initial conditions and amplitude levels A.

The reality of the bounce absorption mechanism is made evident in the phase space plots shown in Fig. 2. In Fig. 2a) one observes the conditions at $\tau = 0$, which correspond to a spatially uniform cold beam travelling toward the wall located at the right edge of the figure. In the absence of the RF field the beam is specularly reflected as indicated in Fig. 2b). However, when the field is present, different particles sample different phases of the oscillation. Thus, a spread in the velocity of the reflected particles occurs and leads to a net absorption of RF power, as shown in Fig. 2c).

It has been verified in the numerical study that the multiple bounces between the walls give rise to a secular increase in the kinetic energy of the charged particle population, as expected. In addition, one finds the existence of a well defined time asymptotic state with a unique value of the kinetic energy per particle, KE/N, for a fixed value of A.

The asymptotic state is independent of the initial phase space configuration, as is demonstrated in Fig. 3. In this figure we exhibit the time evolution of the kinetic energy KE (in arbitrary units) for a fixed $A = 0.2$, but starting from different initial configurations. In this particular example we have considered the cases: (1) slow counter-streaming beams, (2) fast counter-streaming beams, (3) relatively low random velocities with a random spatial distribution, and (4) same as (3) but with fast random velocities. Other exotic initial conditions, including space and velocity bunching, have been considered and show results similar to those of Fig. 3. For sufficiently long times, i.e., $\tau > 50$, all the systems approach a constant value of KE about which fluctuations occur The detailed manner in which the asymptotic state is approached does depend on the initial choice of the configuration; however, as seen in Fig. 3 a well defined ensemble average $\langle KE \rangle$ exists.

The dependence of the ensemble averaged asymptotic kinetic energy per particle $\langle KE \rangle /N$ (in units of $m\omega^2 L$) on the scaled RF amplitude A is shown in Fig. 4. The dots represent individual runs performed with different numbers of particles and/or initial conditions, hence they convey an idea of the statistical nature of the information. The dashed line in Fig. 4 is an eyeball fit to the data which helps to illustrate the overwhelming linear dependence on A over the region $0.1 \leq A \leq 0.5$. In terms of physical variables this dependence implies that

$$\frac{\langle KE \rangle}{N} = \frac{\gamma}{2} e\Phi \tag{13}$$

where $\Phi = 2LE_0$ is the peak RF voltage difference applied between the walls and $\gamma \approx 0.7$ is a constant. This result suggests that the charged particle system attains quasi-equipartition with the externally applied potential independently of the applied frequency, the system size, or the particle mass. The attainment of this pseudo-thermodynamic equilibrium is brought about by the stochastic collisions with the walls.

The dynamical effect which prevents the unbounded growth in KE (i.e., the RF runaway), can be traced to the existence of a maximum velocity $v_M = \omega L/\pi$ for which the direction of the electric field reverses when the particle reaches one of the walls. Once a particle attains a velocity close to v_M no significant energy gain occurs, but such a particle can be scattered into lower velocity bins. Consequently, the system is self-limiting and the maximum energy which can be transferred to it is given by Eq. (13).

IV. Transverse Field Absorption

We proceed to evaluate the power absorbed by the plasma electrons due to a component of the RF electric field which points in a direction $\hat{y}$ parallel to the boundary, as is appropriate for inductive fields, i.e.,

$$\vec{E} = E_0 \exp(-\kappa x) \cos(\omega t)\hat{y} \tag{14}$$

where κ is an evanescent wavenumber which is left unspecified at this stage.

The change in velocity Δv_y experienced by a single particle reflected from the boundary is

$$\Delta v_y = \left(\frac{2eE_0}{m\omega}\right) \frac{(\kappa v_x/\omega)}{1+(\kappa v_x/\omega)^2} \cos\theta \tag{15}$$

where v_x is the speed of the particle in the x direction (unchanged by the boundary) and θ is the random phase of the field at the time the particle arrives at the boundary. The corresponding phase averaged change in the kinetic energy is

$$\langle KE \rangle = \frac{1}{m} \left(\frac{eE_0}{\omega}\right)^2 \frac{(\kappa v_x/\omega)^2}{[1+(\kappa v_x/\omega)^2]} \tag{16}$$

which results in the absorption of RF power per unit area due to all the electrons

$$\frac{P}{A} = \left(\frac{\omega_p}{\omega}\right)^2 \frac{|E_0|^2}{4\pi} \; \bar{v} \;\; G\left(\frac{\kappa\bar{v}}{\omega}\right) \tag{17}$$

where
$$G(z) = \frac{z^2}{\sqrt{\pi}} \int_0^\infty du \; \frac{u^3 \exp(-u^2)}{[1+z^2u^2]^2} \tag{18}$$

for a Maxwellian velocity distribution.

The dependence of the structure function G on $(\kappa\bar{v}/\omega)^2$ is shown in Fig. 5 for the interval $0.01 \leq (\kappa\bar{v}/\omega)^2 \leq 2.01$, i.e., when the transit time is larger than the RF period, and in Fig. 6 for $1.0 \leq (\kappa\bar{v}/\omega)^2 \leq 200$ which corresponds to the small transit time regime. As expected from Eq. (18), in the large transit time limit $G \sim (0.28)(\kappa\bar{v}/\omega)^2$, whereas for small transit times $G \sim (1.08)(\kappa\bar{v}/\omega)^{-2}$. For $(\kappa\bar{v}/\omega) \sim 1$, G exhibits its maximum value, i.e., $G \lesssim 5.64 \times 10^{-2}$.

To express the surface absorption effect in terms of a loading resistance R experienced by the RF generator, one introduces the current I flowing through the conductor, i.e., $E_0 = \omega \ell I$, where ℓ refers to the inductance per unit length of the structure.

Using the planar result of Eq. (17) to approximate the absorption of a cylindrical conductor of radius a and length L yields

$$\frac{R}{L} \equiv \frac{P}{LI^2} = 2\pi a \ell^2 \omega \; \varepsilon_0 \left(\frac{\omega_p}{\omega}\right)^2 \bar{v} \;\; G\left(\frac{\kappa\bar{v}}{\omega}\right) \tag{19}$$

which refers to the resistance per unit length in units of ohms/meter and in which all quantities are evaluated in the MKS system; ε_0 refers to the permittivity of free space. In general, the inductance per unit length is given by $\ell = \mu_0 s$, where s is a unitless quantity involving geometrical factors, hence

$$\frac{R}{L} = 2\,(\mu_0\omega)(k_0 a)(\bar{v}/c)(\omega_p/\omega)^2 s^2\; G(\kappa\bar{v}/\omega) \tag{20}$$

which implies that

$$\frac{R}{|z_0|} = 2\pi\,(k_0 a)(\bar{v}/c)(\omega_p/\omega)^2 s\;\; G(\kappa\bar{v}/\omega) \tag{21}$$

where $|z_0| = \mu_0\omega L s$ is the magnitude of the inductive impedance, μ_0 is the permeability of free space, and $k_0 = \omega/c$ is the vacuum wavenumber.

To obtain a feeling for the possible magnitude of this effect consider the characteristic parameters $\omega/2\pi = 5 \times 10^6$ MHZ, $(\omega_p/\omega)^2 \approx 10^3$, $\bar{v}/c \approx 10^{-2}$, $a \approx 5$ cm, $s \approx 1$, $G \approx 10^{-2}$, which result in $R/|z_0| \approx 3 \times 10^{-2}$ from Eq. (21). This implies that the RF conductor exhibits a quality factor $Q \approx |z_0|/R \approx 30$ due to the plasma boundary absorption alone. If such a RF conductor were to be used to couple energy into collective oscillations of the plasma which have much larger Q values, then it would be found that the loading due to the surface absorption would mask the measurements of the collective mode Q.

The collisionless absorption mechanism described by Eq. (17) has been studied by solid state physicists for many years[7] in the context of the anomalous reflection of microwaves from the surfaces of good conductors, i.e., the anomalous skin effect. Practical applications of the effect have been investigated by Kapitza[8], who has proposed a fusion scheme[9] which derives its electron heating input from this channel. Limited studies of the anomalous skin effect in plasmas have been reported.[10,11]

The connection with the anomalous skin effect studies arises when one demands global conservation of energy. This implies that the evanescent wavenumber in Eq. (14) is quasi self-consistently determined from the

the requirement that the power absorbed, as given by Eq. (17), must be balanced by the Poynting flux into the plasma S_x at $x = 0$, i.e.,

$$S_x = \frac{1}{\mu_0} \, (\vec{E} \otimes \vec{B}) \cdot \hat{x} = \frac{P}{A}$$

$$\frac{\kappa}{\omega\mu_0} \, |E_0|^2 = \frac{P}{A} \tag{22}$$

which yields the consistency relation

$$\frac{\kappa c}{\omega} = \left(\frac{\omega_p}{\omega}\right)^2 \left(\frac{\overline{v}}{c}\right) G \left(\frac{\kappa\overline{v}}{\omega}\right) \tag{23}$$

from which κ is to be determined.

In general, G has the dependence on κ exhibited in Figs. (5) and (6). However, to extract the usually quoted anomalous skin effect one considers only the short transit time regime $(\kappa\overline{v}/\omega) > 1$ and approximates $G \sim (\kappa\overline{v}/\omega)^{-2}$. In this approximate limit Eq. (23) yields

$$\kappa = (\omega_p^2\omega/c^2\overline{v})^{1/3} \tag{24}$$

which gives rise to the absorbed power

$$\frac{P}{A} = \left(\frac{\omega_p}{\omega}\right)^{2/3} \left(\frac{c}{\overline{v}}\right)^{1/3} c \, \frac{|E_0|^2}{4\pi} \tag{25}$$

and the equivalent of Eq. (21) becomes

$$\frac{R}{|z_0|} = 2\pi(k_0 a) \left(\frac{\omega_p}{\omega}\right)^{2/3} \left(\frac{c}{\overline{v}}\right)^{1/3} s \tag{26}$$

The region of validity for Eqs. (25) and (26) is found self-consistently from Eq. (24) to be

$$\left(\frac{\omega_p}{\omega}\right)^{4/3} \left(\frac{\overline{v}}{c}\right)^{4/3} > 1 \tag{27}$$

The collisional RF power absorption P_c which accompanies the anomalous collisionless process is

$$\frac{P_c}{A} = \left(\frac{\omega_p^2}{\omega^2+\nu^2}\right) \left(\frac{c^2\bar{v}}{\omega\omega_p^2}\right)^{1/3} \nu \frac{|E_o|^2}{4\pi} \tag{28}$$

where ν refers to the electron-ion collision frequency.

The ratio of the corresponding collisional loading resistance R_c to the collisionless one is

$$\frac{R_c}{R} = \left(\frac{\omega_p}{\omega}\right)^{2/3} \left(\frac{\bar{v}}{c}\right)^{2/3} \left(\frac{\nu}{\omega}\right) [1+(\nu/\omega)^2]^{-1} \tag{29}$$

which shows that for typical parameters the collisional loading of the RF coupler is at least two orders of magnitude smaller (it can be much smaller) than the collisionless surface absorption loading.

As is well known, in the absence of surface absorption the evanescent wavenumber κ is associated with the collisionless skin depth, i.e., $\kappa = \omega_p/c \equiv \kappa_p$ which arises due to a non-dissipative process. To incorporate this behavior in a continuous manner in the evaluation of κ, which is required in Eq. (17) to obtain the surface power absorption, one can introduce an effective total conductivity σ_T defined by

$$\frac{1}{\sigma_T} = \frac{1}{\sigma_a} + \frac{1}{\sigma_p} \tag{30}$$

where $\sigma_a = \kappa_a^2/i\mu_o\omega$, and κ_a is the anomalous wavenumber given by Eq. (24) or the more exact result obtained from solving Eq. (23); $\sigma_p = i\kappa_p{}^2/\nu_o\omega$ refers to the cold plasma conductivity. The effective wavenumber κ_T associated with σ_T is defined by

$$\kappa_T \equiv (i\mu_o\omega\sigma_T)^{1/2} \tag{31}$$

hence

$$\kappa_T = \left(\frac{\kappa_a^2\,\kappa_p^2}{\kappa_a^2+\kappa_p^2}\right)^{1/2} \tag{32}$$

which shows that as $\kappa_p^2 >> \kappa_a^2$, $\kappa_T \sim \kappa_a$ and for $\kappa_a^2 >> \kappa_p^2$, then $\kappa_T \sim \kappa_p$, as is expected. Physically the $\kappa_p^2 >> \kappa_a^2$ regime refers to low frequency RF applied to a hot plasma and the transition to $\kappa_a^2 >> \kappa_p^2$ implies an increase in frequency and/or a colder plasma. From Eq. (24) the transition is expected to occur at $\omega \sim (\bar{v}/c)\ \omega_p$.

It should be mentioned that under certain conditions a plasma may contain a population of suprathermal electrons[12] whose density n_s is much smaller than the background density, but having an effective thermal velocity v_s much larger than that of the background electrons, i.e., $n_s << n_o$, $v_s >> \bar{v}$. The presence of such electrons can cause an enhancement of the surface power absorption which can be put in the form

$$\frac{P}{A} = \left(\frac{\omega_p}{\omega}\right)^2 \frac{|E_o|^2}{4\pi}\ \bar{v} \left[G\left(\frac{\kappa \bar{v}}{\omega}\right) + \left(\frac{n_s v_s}{n_o \bar{v}}\right) G\left(\frac{\kappa v_s}{\omega}\right) \right] \tag{33}$$

if the velocity distribution of the suprathermal component can be approximated by a Maxwellian. From Eq. (33) it is clear that in certain parameter regimes the background electrons may be found in the $\kappa\bar{v}/\omega < 1$ regime while the suprathermals are in the short transit time limit.

V. Conclusions

The present study provides simple examples which demonstrate that when an unmagnetized plasma surrounds a conductor to which a RF potential is applied, a significant amount of power can be absorbed by the surface electrons due to the process of collisionless boundary absorption. This process consists of both pure bounce absorption and transit time absorption. The surface absorption can lead to an enhanced loading of the RF coupler which in some parameter regimes can mask the loading associated with the far field excitation of collective plasma oscillations. In addition, the process gives rise to energetic surface electrons which may cause a variety of undesirable effects.

It should be noted that the collisionless surface absorption effects are intrinsically kinetic processes which enter into the formal theory as nonlocal terms. Therefore, such effects are absent in coupling calculations which treat the plasma as a cold dielectric, and even in kinetic formulations if the integral equation contributions from the RF boundary are not retained.

A numerical study of pure collisionless bounce absorption heating has demonstrated the physical reality of this process, and has uncovered the existence of a quasi-thermodynamic equilibrium between the charged particles and the applied RF potential.

It is clear that the effect of external magnetic fields on the bounce absorption process is of considerable interest in connection with the RF heating of magnetized plasmas. This is a more difficult problem which deserves investigation.

Acknowledgements

The interest in this topic was motivated by RF coupler measurements[13] performed in collaboration with Dr. R. J. Taylor. Interesting discussions were held with Prof. J. M. Dawson about the anomalous skin effect. This work was supported by the Office of Naval Research.

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FIGURE CAPTIONS

Fig. 1: Frequency dependence of the longitudinal surface power absorption.

Fig. 2: Phase space evolution demonstrating the existence of the bounce absorption process. (a) initial beam configuration, (b) particle reflection without RF, (c) particle reflection with RF present.

Fig. 3: Existence of a time asymptotic state independent of initial conditions: (1) fast counter-streaming beams, (2) slow counter-streaming beams, (3) slow random velocities, random spatial distribution, (4) fast random velocities, random spatial distribution. A = 0.2, total kinetic energy KE is in arbitrary units.

Fig. 4: Dependence of asymptotic kinetic energy per particle (in units of $m\omega^2L/2$) on the RF amplitude $A = eE_0/m\omega^2L$.

Fig. 5: Dependence of the structure function G for transverse surface absorption in the large transit time region.

Fig. 6: Dependence of the structure function G for transverse surface absorption in the small transit time region.

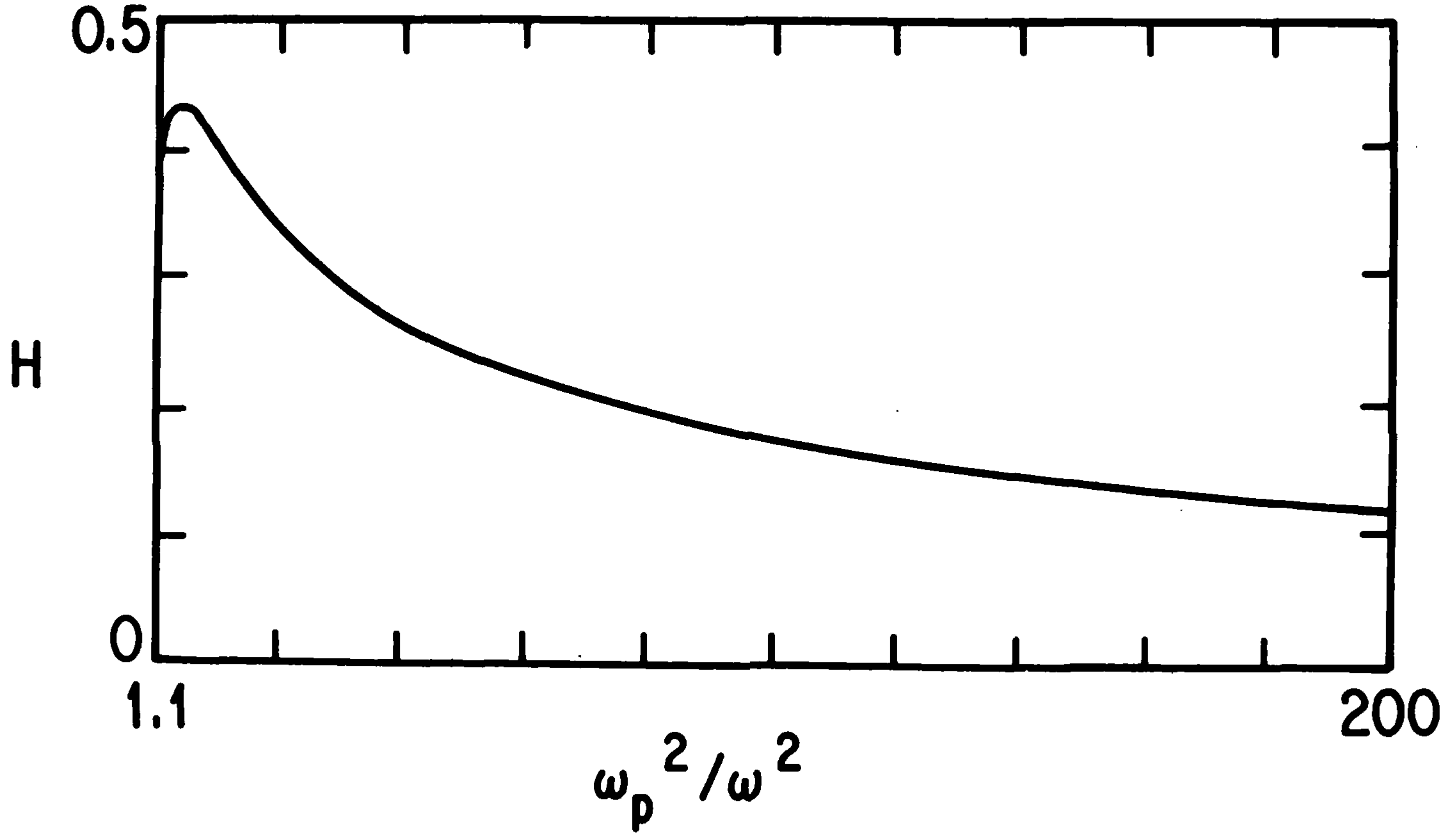

FIG. 1

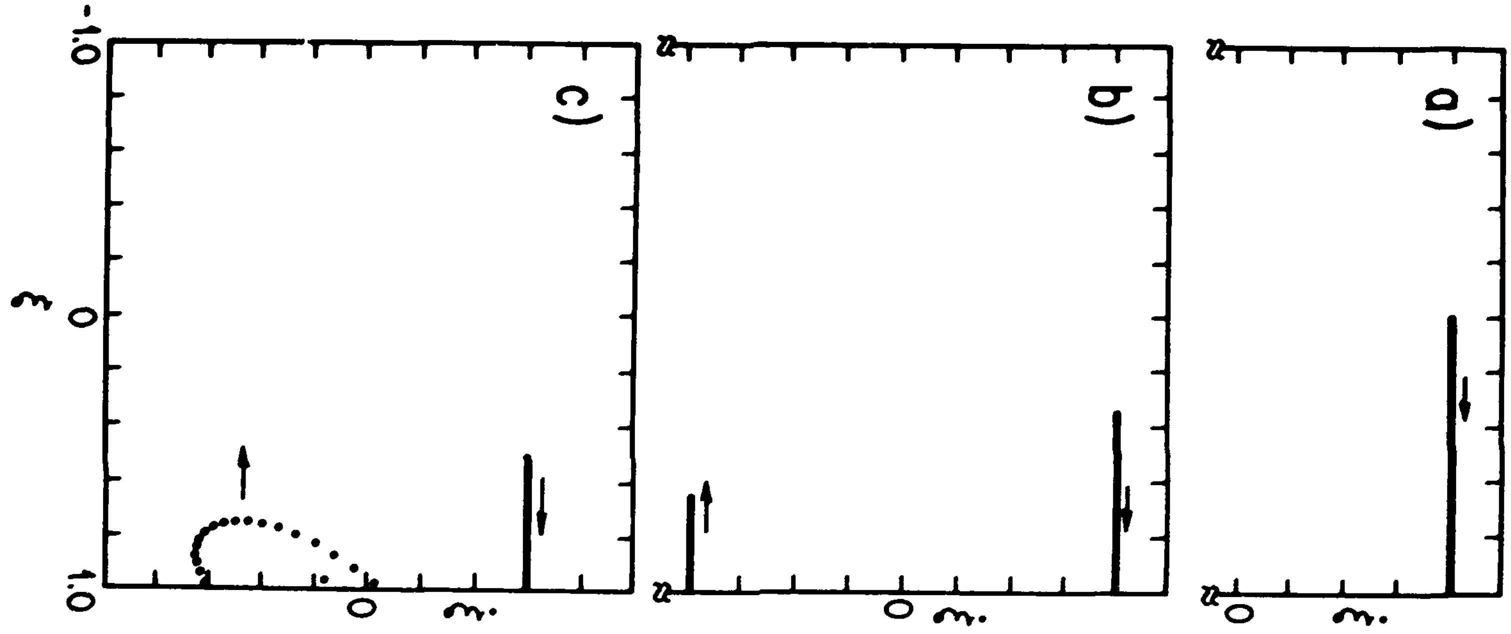

FIG. 2

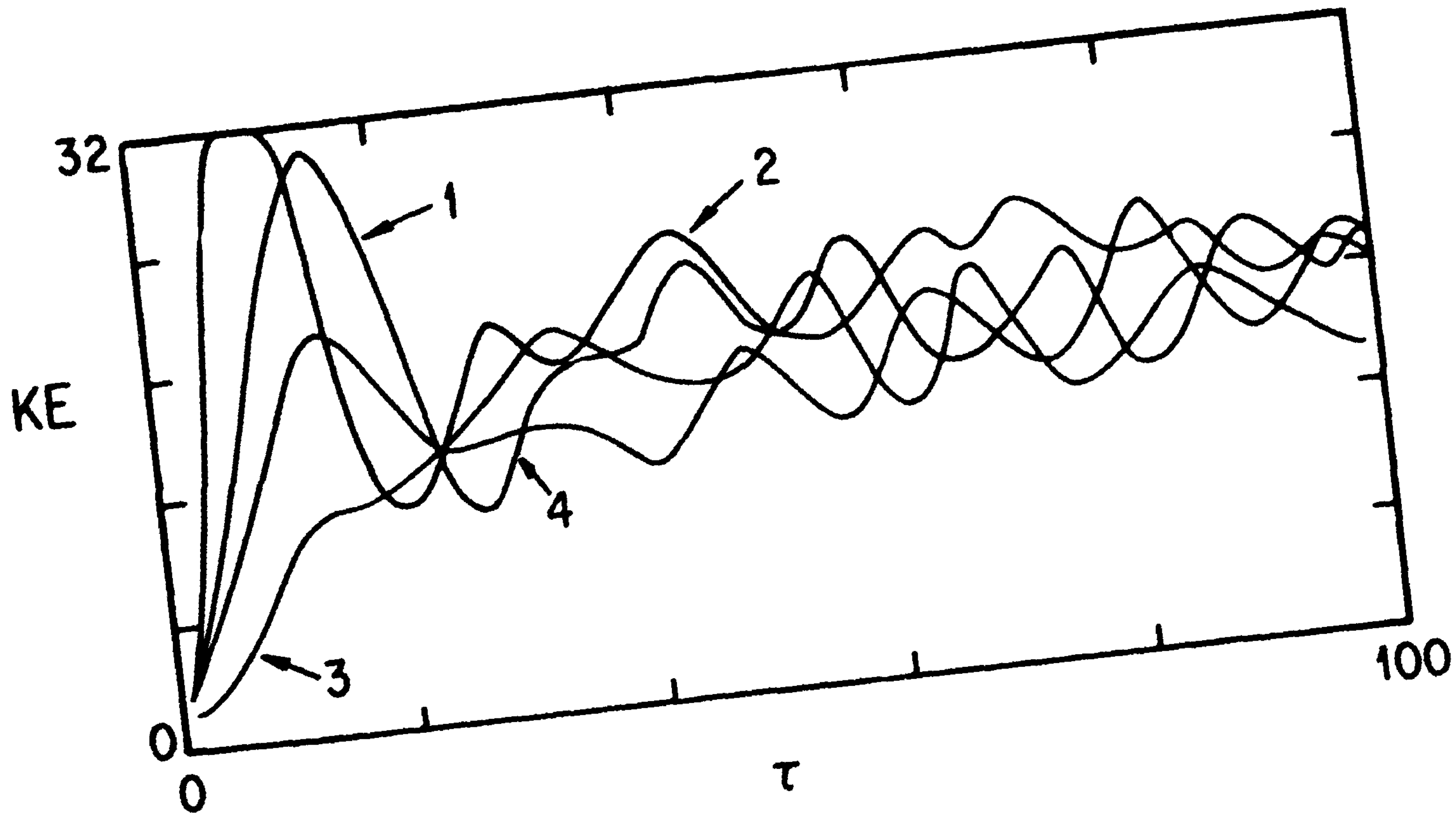

FIG. 3

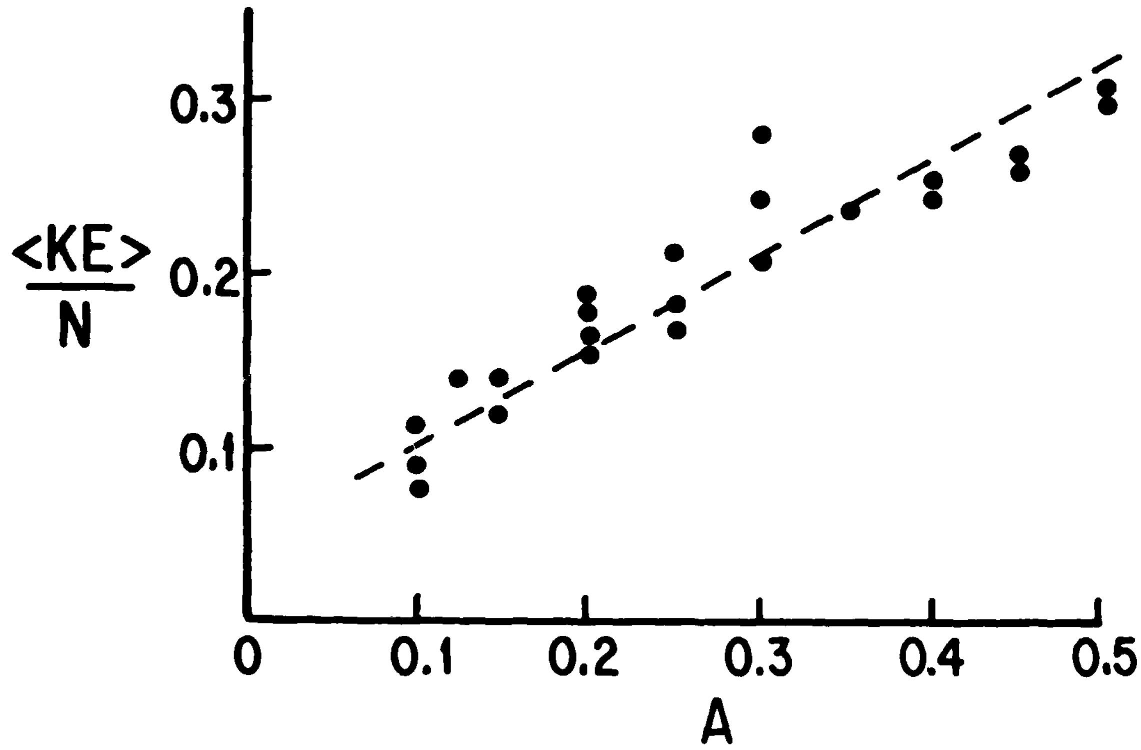

FIG. 4

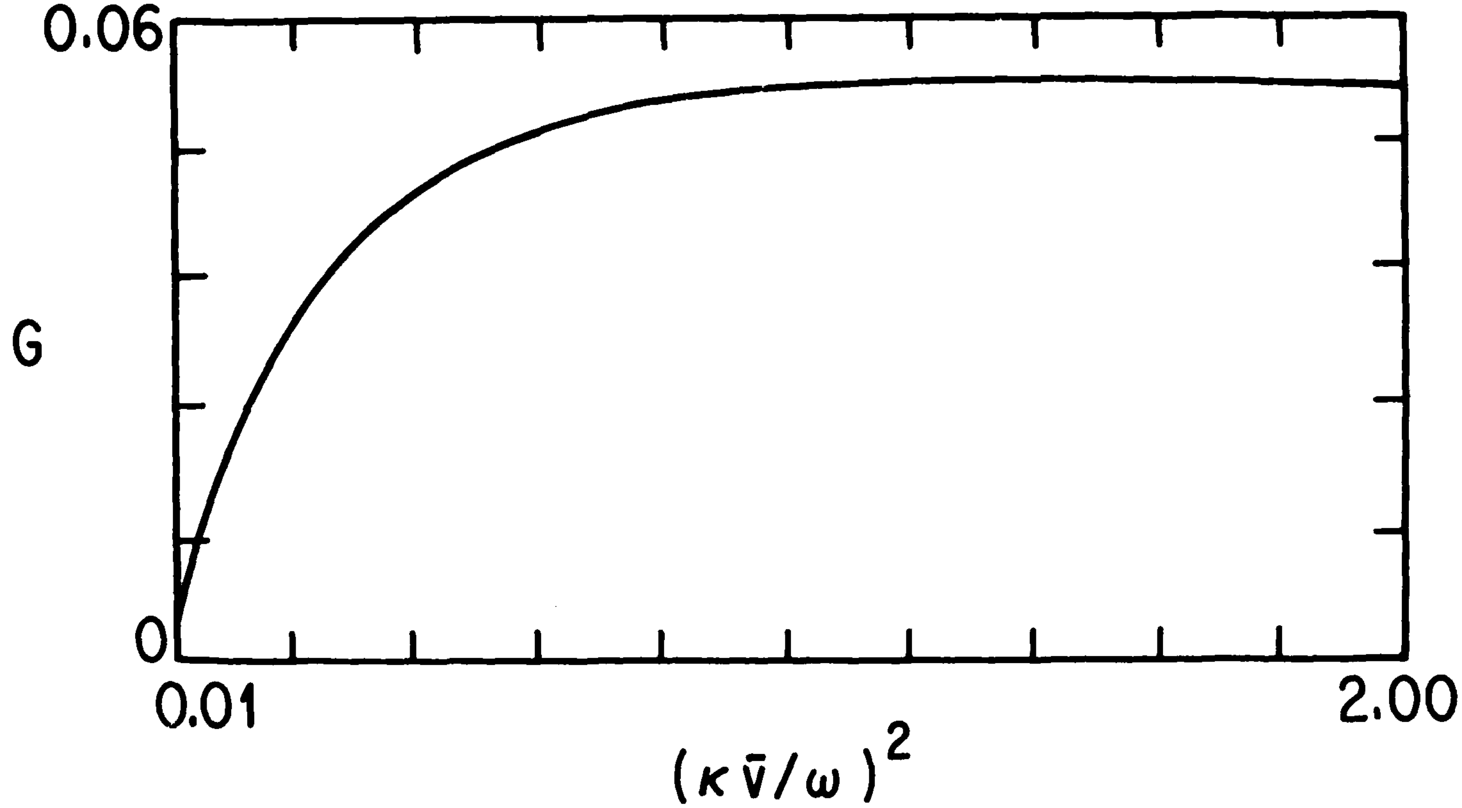

FIG. 5

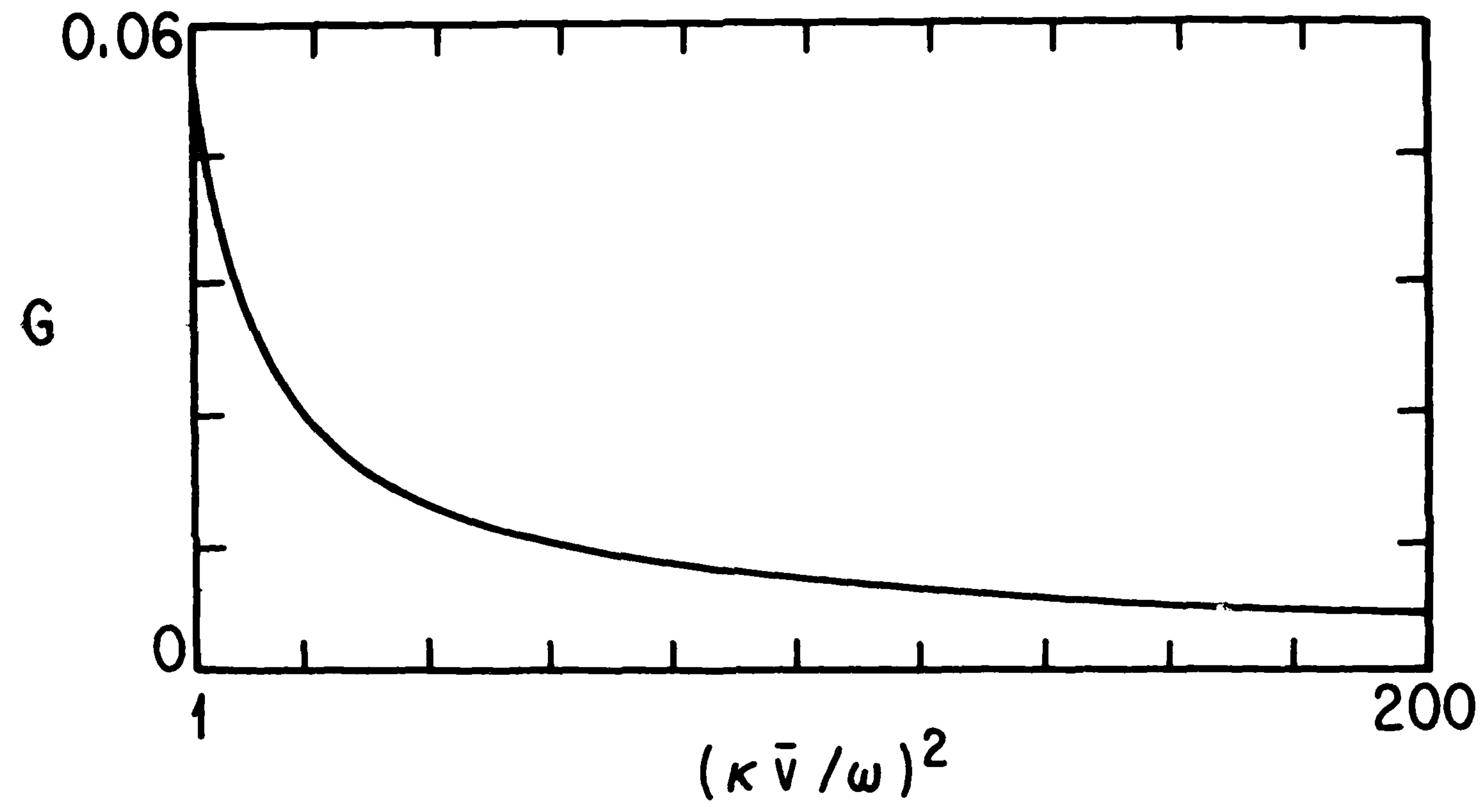

FIG. 6

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